

A low-energy solution to the μ -problem in gauge mediation

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Abstract

In the gauge-mediation framework the soft supersymmetry breaking mass parameters of the supersymmetric standard model are induced by the gauge interactions of some messenger fields. The parameters exhibit flavor universality which is dictated by the gauge interactions and which efficiently eliminates new dangerous contributions to flavor changing neutral currents. However, the Higgs potential in this framework typically contains an unacceptable hierarchy between its dimensionful parameters (the μ -problem of gauge mediation). We show that the problem can be resolved if the Higgs potential arises dynamically once an intermediate $U(1)'$ sector is integrated out rather than arising radiatively from some Yukawa interactions at the messenger scale. As an added benefit, such models may naturally avoid new contribution to CP violating amplitudes. The proposed framework is described, explicit examples are given and its phenomenology is explored. The μ problem is resolved in this case by the low-energy $U(1)'$ dynamics which could be tested in future collider experiments.

I. INTRODUCTION

Supersymmetry provides a well motivated framework for embedding and extending the Standard Model of strong and electroweak interactions (SM): Its boson – fermion symmetry resolves the hierarchy problem and allows one to consistently extrapolate the theory to high-energy. However, a long standing question remains: At what high-energy scale does the next level of structure reside. The answer depends in part on the manner in which supersymmetry breaking is mediated (at a scale $\Lambda_{\text{mediation}}$) from some (hidden) sector in which supersymmetry is broken spontaneously to the SM (observable) sector. It also depends on whether the SM sector itself contains any additional structure, e.g., $\text{SM} \times U(1)'$.

It is possible that the mediation of supersymmetry breaking is carried out through the SM gauge interactions (in which case there may not be a truly hidden sector in the sense that it interacts only gravitationally with the observable sector). Finite gauge quantum corrections involving supermultiplets of postulated heavy (vector-like) matter which transform under the SM – *the messengers* – generate the desired soft supersymmetry breaking (SSB) mass parameters in the low energy potential. One has in this case $m_{\text{SSB}} \sim (\alpha/4\pi)\Lambda$, where Λ is a scale¹ of the order of magnitude of the masses of the messengers (it is given by the ratio of a supersymmetry breaking mass-squared and a supersymmetry conserving mass) and α is a generic gauge coupling. This is the gauge-mediation framework [1]. The physical high-energy scale in this case is set by $\Lambda \sim (4\pi/\alpha)m_{\text{SSB}} \sim (4\pi/\alpha)(1-10) \times M_W \sim 10^{4-5}$ GeV, where m_{SSB} is a typical (SM) superpartner mass. A clear benefit of this framework is the flavor universality of the sfermion \tilde{f} spectrum, which comes about since the different $m_{\tilde{f}}^2 \sim Q_{\tilde{f}}^2[(\alpha/4\pi)\Lambda]^2$ parameters are distinguished only by the respective gauge quantum numbers. This eliminates *a priori* many potentially dangerous sources of flavor changing neutral currents.

The low-mediation scale $\Lambda_{\text{mediation}} \simeq \Lambda \sim 10^{4-5}$ GeV renders negligible, in most cases, any supergravity effects and hence eliminates various supergravity schemes for the generation of the μ -parameter – the supersymmetry conserving Dirac mass which mixes the two Higgs doublets of the supersymmetric extension (SSM) $W = \mu H_1 H_2 + \dots$. Hence, the μ problem, why $\mu \simeq \mathcal{O}(M_W)$ rather than $\mu \simeq \mathcal{O}(\Lambda_{\text{mediation}})$, resurfaces in this case. (The non-observation of chargino pairs at the WW threshold implies that $|\mu| \gtrsim M_W$, and in particular, that it cannot vanish [2].) One can consider the possibility that the dimension-one μ -parameter is generated by Yukawa quantum corrections which involve some messenger fields which interact with the Higgs doublets of the SSM via Yukawa couplings y . However, if this were the case then it is straightforward to show [3] that a dimension-two SSB mixing between H_1 and H_2 in the scalar potential $V_{\text{SSB}} \sim \dots + m_3^2 H_1 H_2 + h.c. + \dots$ would also be generated and at the same loop order. Unlike in the case of gauge loops, dimension one and two parameters generated by Yukawa loops are typically suppressed by the same power of the loop factor. As a result $|\mu| \sim (y^2/16\pi^2)\Lambda \sim m_3^2/\Lambda \sim M_W$ (or equivalently, $\mu^2 \sim (16\pi^2/y^2)m_3^2$). It reintroduces a hierarchy problem to the Higgs potential which would be dominated in this

¹ The messenger scale may be roughly the same as or smaller by a few orders of magnitude than the actual scale of supersymmetry breaking.

case by $m_3^2 \sim M_W \Lambda$. This is a most severe problem that undermines any success of the gauge-mediation framework.

The most successful attempts to address this new hierarchy problem involve in one fashion or another the details of the high-energy (supergravity) theory, and in that sense they are high-energy solutions. For example, Ref. [3] invokes a radiative linear term generated by messenger-scale singlet interactions. The linear term shifts a singlet field N (which interacts with the Higgs doublets) to a scale which is suppressed by a loop factor in comparison to the messenger scale. The shifts in the scalar and auxiliary components of N , which induce μ and m_3^2 respectively, arise at different loop orders, evading the above described hierarchy problem. The superpotential (or equivalently – the Kähler potential) couplings must be fixed by the high-energy ($Q > \Lambda$) theory. In particular, a scale associated with a tree-level linear term must be fixed to be $\mathcal{O}(\Lambda)$. Alternatively, in Ref. [4] it was pointed out that a radiative linear term in a singlet field N is typically generated by supergravity and is suppressed by only one inverse power of the Planck mass M_P . Hence, it can still play an important role in the low-energy theory. It shifts the singlet field $N \sim (\Lambda^4/\kappa^2 M_P)^{1/3}$ (assuming $W \sim (\kappa/3)N^3$). The singlet Yukawa interaction with the Higgs doublets then generates the desired parameters at tree-level $\mu^2 \sim m_3^2 \sim N^2 \sim M_W^2$ (assuming that supersymmetry is broken at a scale of the order of $10^{6\pm 1}$ GeV). In this case no new scales are introduced by hand, but there is still dependence on the high-energy theory. (A somewhat similar application of supergravity to the problem was proposed in Ref. [5].) Both solutions assume the presence of a “dedicated” messenger-scale gauge singlet(s), denoted above by N .

Here, we point out a distinctive possibility that the singlet field is not a gauge singlet but only a SM singlet S . Specifically, we assume the extension $(S)SM \rightarrow (S)SM \times U(1)'$, and that S carries a charge $Q_S = -(Q_{H_1} + Q_{H_2})$ under the additional Abelian symmetry so that a Yukawa term $W \sim h_s S H_1 H_2$ is allowed. In turn, a scale $\Lambda' \sim \langle S \rangle \lesssim \Lambda$, which is associated with the breaking of the $U(1)'$, must be introduced, or preferably, induced. The various μ problems of gauge mediation will be shown to be solved in this case by the low-energy dynamics associated with this new scale.

The scale Λ' could be generated radiatively and is a function in this case of Λ and of $\mathcal{O}(1)$ Yukawa couplings. A coupling between S and exotic quarks, e.g., D and D^c singlets with hypercharge $\pm(1/3)$, generates negative corrections to m_S^2 so that $m_S^2(\Lambda') < 0$ and S acquires a vacuum expectation value (vev). This is essentially a $U(1)'$ version of the well-known radiative symmetry breaking (RSB) mechanism that is responsible in the SSM for the generation of the negative mass term in the SM Higgs potential. A similar idea was suggested previously in the context of supergravity and high-energy (gravity) mediation ($\Lambda_{\text{mediation}} \simeq M_P$) of supersymmetry breaking [6]. In that case, like RSB in those models, the large evolution interval enables one to render $m_S^2 < 0$ somewhere above the weak scale. In the supergravity case the superpotential interactions generate $|\mu| \sim h_s \langle S \rangle$ while trilinear SSB terms $V_{SSB} \sim \dots + h_s A_s S H_1 H_2 + h.c. + \dots$ generate $m_3^2 = A_s h_s \langle S \rangle$. Since all parameters in the gravity-mediation framework are of the same order of magnitude as the gravitino mass (which is fixed $m_{3/2} \sim M_W$), then $h_s \langle S \rangle$ is expected to be of the same order of magnitude as well. This leads to a successful solution to the μ -problem in high-energy supergravity models. Various other solutions that benefit from the large mediation scale are also available in the supergravity framework [7]. The radiatively broken $U(1)'$ scenario was extensively

studied in the framework of supergravity and of string-inspired models [8].

In contrast to the supergravity framework, in gauge mediation the evolution interval is short; in addition, trilinear parameters are highly suppressed² $A \sim (\alpha/4\pi)^2 \Lambda \ln \Lambda$. While the small A parameters remain a constraint, the shorter evolution interval is more than compensated (as for the case of RSB in these models) by the large hierarchy within the SSB parameters $m_D^2/m_H^2/m_S^2 \sim \alpha_3^2/\alpha_2^2/\alpha_1^2$ (where $\alpha_{3,2,1,1'}$ are the $SU(3)$, $SU(2)$, $U(1)$, and $U(1)'$ gauge couplings). In fact, the messengers may not transform under $U(1)'$, in which case $m_S^2(\Lambda) = 0$. For $\alpha_{1'} = \mathcal{O}(\alpha_Y)$, which we will assume, the exact boundary condition for m_S^2 does not affect our discussion and for simplicity we assume hereafter that the messengers are indeed invariant under $U(1)'$.

A radiatively induced $\langle S \rangle$ as a source of μ in the case of a gauge singlet S was considered previously in the context of gauge mediation [10]. It was found that the singlet must couple to exotic quarks with large Yukawa couplings, as naturally occurs in the context of $U(1)'$. In the gauge singlet case, however, the superpotential must contain a S^3 term so that the potential contains quartic terms $V \sim |\partial(S^3 + SH_1H_2)/\partial S|^2$ which stabilize it. The models suffer from the usual problem of a spontaneously broken global Z_3 symmetry (under which S^3 is invariant) which results in unacceptable domain walls at a (post-inflationary) low-energy epoch. In the gauged case S is not a singlet and S^3 terms are not gauge invariant and are automatically forbidden. Instead, the potential is stabilized by $U(1)'$ gauge D -terms $V \sim \dots + (g_1^2/2)(Q_S|S|^2 + Q_{H_1}|H_1|^2 + Q_{H_2}|H_2|^2)^2 + \dots$ (which are not available for a gauge singlet S). The Z_3 symmetry is now only a (harmless) subgroup of the gauged $U(1)'$. While in the non-gauged case the former source of the quartic terms also generates an additional contribution to $m_3^2 \sim S^2$, this is not possible in the gauged case (with only one singlet).

In either the gauged or non-gauged case, the potential also exhibits an approximate phase (R) symmetry, which exists in models with only Yukawa superpotential terms and corresponds to a rotation of all fields by the same phase. It is broken spontaneously by $\langle S \rangle$ and explicitly by tri-linear A -terms. The explicit breaking is, however, suppressed by the smallness of the A -parameters. We will show below that in spite of the suppressed A -parameters it is possible to generate m_3^2 and break the phase symmetry strongly enough to avoid the light pseudo Goldstone boson which otherwise appears. Specifically, as will be shown below, it is very likely that in the $U(1)'$ scenario $M_W \ll \langle S \rangle \lesssim \Lambda$ and hence, $m_3^2 \sim h_s A_s \langle S \rangle \sim A_s \mu$ is a geometric mean of a small parameter and a large vev. It implies a somewhat large value of $|\mu| \sim \mathcal{O}(1 \text{ TeV})$. However, this typically occurs in gauge mediation as a result of RSB constraints in the presence of a heavy gluino [11]. The details depend strongly on the exotic quark spectrum and on the $U(1)'$ charges, and some examples will be presented below. Alternatively, in models with two singlets a superpotential term SS'^2 could be gauge invariant, and $\langle S' \rangle \sim \langle S \rangle$ could generate an additional contribution to $m_3^2 \sim \langle S' \rangle^2$, just as in the non-gauged case. (Note that in the non-gauged case the $U(1)'$ rotations – explicitly broken by the S^3 terms – correspond to global transformations and there is one additional pseudo Goldstone boson.)

²Large A parameters were proposed, however, in Ref. [9].

It is particularly interesting to note that in the models with only one SM singlet there appear only two new phases which can be rotated away, and hence there are no new physical phases³. This is because there is only one common phase to all gaugino mass and the radiatively-induced A parameters, while the phase of m_3^2 is given in this case by the phases of μ and A . Hence, after R and Peccei-Quinn rotations no physical phases appear in the soft parameters. This eliminates new contributions to CP violating amplitudes such as the electron dipole moment, which are flavor conserving and which generically appear at unacceptable levels even in gauge-mediation models [12].

The stabilization due to the D -terms and the generation of the A terms then open the door to new (low-energy) solutions to the μ -problem in gauge mediation. The mechanism is quite different from that of the non-gauged case since the quartic coupling is given, in principle, by a fixed gauge coupling rather than by a free superpotential coupling; m_3 must depend on overcoming the suppression of the tri-linear couplings A ; and the scale $\langle S \rangle$ is a physical scale with observable consequences. Hence, it corresponds to a distinctive and interesting option and it will be explored in some detail below.

The $U(1)'$ models predict, in addition to the extra matter and the associated rich spectrum, an extra gauge boson, Z' . The corresponding phenomenology is similar to that of any other model with Z' , except that $M_{Z'}^2 \sim -(Q_S/2)m_S^2$ is large, given that $|m_S^2|$ is controlled by the large exotic quark SSB parameters. Typically we find $M_{Z'} \simeq \mathcal{O}(1 \text{ TeV})$ and with suppressed mixing with the ordinary Z -boson. Thus it decouples safely from electroweak physics [13]. Another interesting aspect of supersymmetric $U(1)'$ models that repeats here is that the tree-level light Higgs h_1 mass exceeds its usual upper bound of M_Z . This is due to contributions from the $U(1)'$ D -terms to the quartic potential, which lift its otherwise flat direction. We find for its mass $m_{h_1} \simeq 120 - 150 \text{ GeV}$ at tree level and $m_{h_1} \simeq 150 - 180 \text{ GeV}$ at one loop.

A most interesting aspect of the $U(1)'$ scenario is that the gauge-mediation scale is still the only fundamental scale, and the $U(1)'$ scale is determined from it. It has been proposed recently [14] that perhaps the same $U(1)'$ is also responsible for the actual mediation of supersymmetry breaking from the “hidden” sector to the messenger fields (i.e., to identify $U(1)'$ with $U(1)_{\text{messenger}}$ of Ref. [1]). This is an ambitious yet interesting proposal that significantly differs from our bottom-up approach, which, in principle, is independent of the details of supersymmetry breaking and its initial mediation to the messenger fields. By distinguishing the two extended interactions we avoid the need, e.g., to fine tune Yukawa couplings, which is the situation in Ref. [14] due to the multitude of tasks imposed there on a single $U(1)$. The only (moderate) hierarchy in Yukawa couplings that is assumed is between those that involve (exotic) quarks, which are taken to saturate or be near their infra-red quasi-fixed points and be $\mathcal{O}(1)$, and those which involve only the Higgs doublets and the singlet(s), which do not reach any (quasi-)fixed points and hence are taken to be smaller. Such differences naturally stem from QCD renormalization, which enables the existence of quasi-fixed points for the (exotic) quark couplings.

³ We thank J. Feng and T. Moroi for bringing the CP problem in gauge mediation and its possible solution to our attention.

We describe the model in greater detail in the next section. In Section III we discuss the evolution of its parameters, the scalar potential, and the extraction of μ and m_3^2 in the two schemes. We summarize and comment on various issues such as unification in the last section. Since our focus in this work is on the demonstration of the viability of the proposed low-energy solution to the μ -problem in gauge mediation, some assumptions (which are clearly indicated) were made in order to simplify the presentation and discussion (mainly by fixing parameters whose variation was found to be immaterial for our purposes). Also, only results for the Higgs and neutralino/chargino sectors which significantly differ from the usual case will be presented in detail, though we checked that the complete spectrum is consistent. Our presentation is within a one-scale model though generalizations are straightforward. Our focus is on the low-energy theory below the messenger scale and on the generation of the $U(1)'$ and electroweak scales as functions of this scale. We comment on implications for possible extrapolations to higher energies in our conclusions.

II. THE MODEL

The extended gauge symmetry of the model is $SU(3)_c \times SU(2)_L \times U(1)_Y \times U(1)'$, with the gauge couplings g_3 , g_2 , $g_1 = \sqrt{\frac{5}{3}}g_Y$ and $g_{1'}$, respectively, and $\alpha_i = g_i^2/4\pi$. The extended particle content is given by the left-handed chiral multiplets of the the minimal supersymmetric extension (MSSM) including three families of quark doublets (Q) and singlets (u^c and d^c), lepton doublets (L) and singlets (e^c), and the Higgs doublets H_1 and H_2 . In addition, it contains exotic quark vector-like pairs, e.g. D and D^c pairs which are singlets under $SU(2)_L$ and carry hypercharge $\pm 1/3$, and fields which are singlets under the SM but are charged under $U(1)'$. We assume that all of the low-energy matter fields (but not the messenger-scale fields) are charged under $U(1)'$. For concreteness, we further assume in most examples that the $U(1)'$ charges of the fields are given by the $U(1)_\eta$ of the E_6 model [15], up to an overall normalization. The E_6 symmetry is used as a classification tool only, i.e., we do not assume a full grand unification. Nevertheless, it enables one, in principle, to choose an anomaly free low-energy spectrum. In this study we include only a subset of fields which are relevant for the generation of the μ and m_3^2 parameters and only comment on the anomaly cancellation in our conclusions. Also, in order to explore a wider range of possibilities, the $U(1)_\eta$ assumption will be relaxed in certain cases, and we will vary the $U(1)'$ charge assignments.

The superpotential of the model is

$$W = h_s S H_1 H_2 + h_u u_3^c Q_3 H_2 + h_d d_3^c Q_3 H_1 + h_e e_3^c L_3 H_1 + h_D S D_i D_i^c + (\text{self-couplings of } S, S') \quad (1)$$

In Eqn. (1) we include the usual Yukawa terms involving the third generation fields, an effective μ term $h_s S H_1 H_2$, and indicate possible self-couplings of the singlet fields.

The soft supersymmetry breaking parameters are generated at the messenger scale Λ through gauge mediation. The breaking of the supersymmetry is parameterized by the singlet X , whose scalar component acquires a vev $\langle X \rangle$ and auxiliary component F a vev $\langle F_X \rangle$. The messenger fields Φ and $\bar{\Phi}$ are vector-like pairs which transform under the SM

gauge group. Here, we assume only one such pair of $\mathbf{5}$ and $\bar{\mathbf{5}}$ of $SU(5)$. The superpotential term

$$W = \lambda X \Phi \bar{\Phi}, \quad (2)$$

where λ is a Yukawa coupling, generates the supersymmetry conserving and breaking masses for the messenger fields $\sim \lambda \langle X \rangle$ and $\sqrt{|\lambda \langle F_X \rangle|}$, respectively.

Because Φ and $\bar{\Phi}$ are charged under the SM gauge group, the effect of the supersymmetry breaking is propagated at the quantum level to the observable sector via the usual gauge interactions. The messenger scale is defined here as $\Lambda = \langle F_X \rangle / \langle X \rangle$. The gaugino masses are generated at one-loop [1]

$$M_i = \frac{\alpha_i}{4\pi} r_i \Lambda, \quad (3)$$

where $i = 1', 1, 2, 3$, and r_i is the Dynkin index for Φ and $\bar{\Phi}$. We choose, as is customary, a normalization such that the minimal $\mathbf{5} + \bar{\mathbf{5}}$ model has $r_1 = r_2 = r_3 = 1$. In particular, we assume⁴ (for simplicity only) that the messenger fields are not charged under the additional $U(1)$, and hence, $r_{1'} = 0$. The scalar masses arise at the two-loop level and are given by

$$m^2 = 2\Lambda^2 \sum_i \left(\frac{\alpha_i}{4\pi}\right)^2 C_i, \quad (4)$$

where C_i are the quadratic Casimirs of the observable sector gauge groups, i.e., $C_3 = 4/3$ for $SU(3)_c$ triplets, $C_2 = 3/4$ for $SU(2)_L$ doublets, and $C_1 = \frac{3}{5}Q_Y^2$ for the hypercharge. Again, the $U(1)'$ does not contribute to the m^2 parameters since by assumption Φ and $\bar{\Phi}$ have zero $U(1)'$ charges. We will take eqs. (3) and (4) to be the boundary condition at a scale Λ , though more generally one could choose a slightly different scale for the boundary. (Our results do not depend on this assumption.)

The A parameters arise only at the higher-loop order and are very small. In our study we take the A parameters to be zero at the messenger scale. Nevertheless, non-trivial values of the A parameters arise from the one-loop renormalization-group evolution. Assuming that the messenger fields Φ and $\bar{\Phi}$ are not charged under the $U(1)'$, the gaugino of the $U(1)'$ and the SM singlet fields are therefore massless at the messenger scale⁵. However, $m_{\tilde{S}}^2$ acquires a non-zero value at the low energy scale due to loop corrections. In particular, it is driven

⁴ This assumption is formally inconsistent with an E_6 embedding of the messengers $\mathbf{5} + \bar{\mathbf{5}} \subset \mathbf{27}$, and is therefore inconsistent with a true E_6 embedding of the models, which we do not assume at this point.

⁵ Even if Φ and $\bar{\Phi}$ are charged under the $U(1)'$, it will not affect the conclusions we have reached here, since the contributions from the $U(1)'$ to the soft mass parameters are, in general, small. It can affect, however, the (singlet) slepton spectrum, which is otherwise given only by hypercharge loops.

rapidly to large and negative values due to the large couplings between S and the exotic quark pairs D and D^c . Once $U(1)'$ is broken, its gaugino and gauge boson are degenerate in mass.

We use the renormalization group equations (RGE) to relate the boundary conditions for the SSB parameters at the messenger scale to their values at lower energies. As the first step, the scale at which the additional $U(1)'$ is broken is determined and is required to be higher than the electroweak scale (in order to be consistent with the experimental limits on the Z' gauge boson). At this scale, the exotic quarks acquire heavy masses and decouple from the theory. By iterating this procedure, the gauge couplings at the messenger scale $\Lambda = 10^5$ GeV are determined from their well-known electroweak-scale values (taking into account the contributions of the exotic matter). At the second step, the values of μ (or equivalently, h_s) and $\tan\beta = \langle H_2 \rangle / \langle H_1 \rangle$ at the minimum of the Higgs potential are determined, the former by fixing the mass of the Z boson and the latter by using the result for μ and $m_3^2 = A_s \mu$. Given $\tan\beta$, the Yukawa couplings at the electroweak scale are determined from the respective fermion masses. The procedure is then iterated in order to determine the correct Yukawa couplings for the t and b quarks and the τ -lepton. The desired solution for the minimum of the Higgs potential which correctly reproduces the Z , t , b and τ masses is found and the Z' mass and mixing, in addition to the sparticle spectrum, are predicted. Given our assumptions, the free parameters in the analysis are Λ , the number n_D of D , D^c pairs that couple to S , the corresponding Yukawa couplings h_D , and (in the case of two SM singlets) the singlet self-coupling. The product of n_D and h_D is constrained by electroweak breaking and also by requiring a sufficiently heavy Z' . Below we fix $n_D = 3$ and use different values of $h_D = \mathcal{O}(1)$. We comment on the variation of the parameters at the end of the next section.

III. THE RADIATIVELY INDUCED HIGGS MASS PARAMETERS

The Higgs potential contains three contributions

$$V = V_F + V_D + V_{soft}. \quad (5)$$

In the following two subsections, we study two cases with one or two SM singlet fields in the model. In the second case, the additional singlet S' has a coupling SS'^2 allowed by gauge invariance. The most general renormalizable superpotential involving the Higgs and two singlet fields is

$$W = h_s S H_1 H_2 + h_{s'} S S'^2. \quad (6)$$

The potential (5) is given in this case by

$$V_F = |h_s H_1 H_2 + h_{s'} S'^2|^2 + |h_s S|^2 (|H_1|^2 + |H_2|^2) + 4|h_{s'} S S'|^2; \quad (7)$$

$$V_D = \frac{G^2}{8} (|H_1|^2 - |H_2|^2)^2 + \frac{g_2^2}{2} |H_1^\dagger H_2|^2 + \frac{g_{1'}^2}{2} (Q_1 |H_1|^2 + Q_2 |H_2|^2 + Q_S |S|^2 + Q_{S'} |S'|^2)^2; \quad (8)$$

$$V_{soft} = m_1^2|H_1|^2 + m_2^2|H_2|^2 + m_S^2|S|^2 + m_{S'}^2|S'|^2 + (A_s h_s S H_1 H_2 + \text{h.c.}) + (A_{s'} h_{s'} S S'^2 + \text{h.c.}), \quad (9)$$

where $G^2 = g_Y^2 + g_2^2$. The one singlet case is given by simply setting $h_{s'} = 0$ and $Q_{S'} = 0$.

A. Radiatively breaking the $U(1)'$ symmetry

The experimental constraint on the mass of the Z' can be satisfied if the $U(1)'$ is broken at the TeV scale, which requires $\langle H_1 \rangle, \langle H_2 \rangle \ll \langle S \rangle, \langle S' \rangle$. This separation is indeed realized in our examples and the determination of the S vev can therefore be separated to a very good approximation from that of the Higgs vev's. We will illustrate the radiative breaking of the $U(1)'$ symmetry in the case of a single SM singlet S . The scalar potential for S reads

$$V = m_S^2|S|^2 + \frac{g_{1'}^2}{2}(Q_S|S|^2)^2. \quad (10)$$

It acquires a vev $\langle S \rangle = s/\sqrt{2}$ where

$$s^2 = -\frac{2m_S^2}{g_{1'}^2 Q_S^2}, \quad (11)$$

if the evolution of m_S^2 can be neglected near the minimum. Hence, a large value for s occurs for m_S^2 large and negative. This is achieved by the order unity Yukawa couplings between S and exotic quark pairs D and D^c (with scalar mass-squares $m_{D, D^c}^2(\Lambda) \gg m_S^2(\Lambda) \simeq 0$), which rapidly diminish $m_S^2(Q < \Lambda)$ via the usual renormalization group evolution. The mass of the Z' boson, which is independent of $g_{1'}$, is

$$M_{Z'} \sim g_{1'} Q_S s \sim \sqrt{2|m_S^2|}, \quad (12)$$

with the $Z - Z'$ mixing angle $\alpha_{Z-Z'} = \mathcal{O}(M_Z^2/M_{Z'}^2)$. The Z' mass and the $U(1)'$ scale are determined by the only scale in the problem, Λ (which is encoded in m_S^2).

B. One singlet models

We first consider models with only one singlet field, with its $U(1)'$ charge satisfying $Q_S + Q_{H_1} + Q_{H_2} = 0$. The superpotential is given by the first five terms in (1). The vev of S generates an effective μ parameter $\mu = h_s s/\sqrt{2}$. The A -term associated with $S H_1 H_2$, which is non-zero at the electroweak scale due to loop corrections, generates an effective m_3^2 for the two Higgs doublets $m_3^2 = A_s \mu$. In addition, the $U(1)'$ D -term generates corrections to the Higgs scalar masses $\delta m_{1,2}^2 = \frac{g_{1'}^2}{2} Q_{1,2} Q_S s^2$. Defining $\langle H_1^0 \rangle = v_1/\sqrt{2}$ and $\langle H_2^0 \rangle = v_2/\sqrt{2}$, the Higgs potential for $v_{1,2}$ at the minimum for s is

	Λ	M_Z		Λ	M_Z
h_u	0.84	0.98	h_d	0.30	0.42
h_e	0.17	0.18	h_D	0.70	0.84
h_s	0.47	0.40	A_s (GeV)	0	-53
A_u (GeV)	0	449	A_d (GeV)	0	573
A_e (GeV)	0	24	A_D (GeV)	0	494
m_1^2 (GeV) ²	(350) ²	(83) ²	m_2^2 (GeV) ²	(350) ²	-(778) ²
m_S^2 (GeV) ²	0	-(821) ²	$m_{D(c)}^2$ (GeV) ²	(1310) ²	(1440) ²

TABLE I. The initial and final values for the first example of the one singlet case. $n_D = 3$ and $\Lambda = 10^5$ GeV.

$$V = \frac{\mu^2}{2}(v_1^2 + v_2^2) + \frac{G^2}{32}(v_2^2 - v_1^2)^2 + \frac{1}{2}(m_1^2 + \delta m_1^2)v_1^2 + \frac{1}{2}(m_2^2 + \delta m_2^2)v_2^2 - m_3^2 v_1 v_2 + \dots, \quad (13)$$

where we use a suitable gauge rotation to render v_1 and v_2 real and positive, and we neglect small corrections from the $U(1)'$ D term which are quadratic in v_1/v_2 . (The usual MSSM loop corrections can also be absorbed in δm_i^2 , but are a secondary effect at this level.) We present in the following two numerical examples with different choices of Q_1 and Q_2 .

In the first example, we choose the $U(1)_\eta$ assignments $Q_1 = 1$, $Q_2 = 4$ and $Q_S = -Q_1 - Q_2$. The initial and final values of the parameters are listed in Table I. We take $\Lambda = 10^5$ GeV and 3 pairs of exotic quark singlets.

The vev of the singlet is $s = 3720$ GeV, the vev's of the Higgs doublets are $v_1 = 14$ GeV and $v_2 = 245$ GeV, resulting in a solution with $\mu = 1050$ GeV and $\tan \beta = 18$. The effective m_3^2 is $\sim (235 \text{ GeV})^2$. The Z' mass is $M_{Z'} = 1110$ GeV and the $Z - Z'$ mixing angle is $\alpha_{Z-Z'} = 0.004$. The (tree-level) spectrum of the CP even physical Higgs is $m_{h_1} = 124$ GeV, $m_{h_2} = 995$ GeV, $m_{h_3} = 1090$ GeV, while $m_{h_1} = 154$ GeV at one loop (with negligible corrections to $m_{h_{2,3}}$). The CP odd Higgs scalar and the charged Higgs masses are $m_A \simeq m_{H^\pm} = 993$ GeV. The heaviest CP even Higgs scalar h_3 is mainly composed of the singlet S , associated with the breaking of the $U(1)'$. The second heaviest CP even Higgs, the CP odd Higgs and the charged Higgs fields form the $SU(2)$ doublet that is not associated with the $SU(2) \times U(1)_Y$ breaking.

The masses of the two charginos are $m_{\tilde{\chi}_1^\pm} = 266$ GeV and $m_{\tilde{\chi}_2^\pm} = 1060$ GeV. The lightest (heaviest) chargino is predominantly a gaugino (Higgsino). The spectrum of the neutralinos is $m_{\tilde{\chi}_1^0} = 142$ GeV, $m_{\tilde{\chi}_2^0} = 266$ GeV, $m_{\tilde{\chi}_3^0} = 1060$ GeV, $m_{\tilde{\chi}_4^0} = 1060$ GeV, $m_{\tilde{\chi}_5^0} = 1120$ GeV, $m_{\tilde{\chi}_6^0} = 1120$ GeV. In the limit of neglecting v_1 and v_2 , the two lightest neutralinos are just \tilde{B} and \tilde{W}_3 , i.e., the Bino and the Wino. $\tilde{\chi}_{3,4}$ are linear combinations of Higgsinos with nearly degenerate masses $\sim \mu = h_s s / \sqrt{2}$; and $\tilde{\chi}_{5,6}$ are linear combinations of the other gaugino \tilde{B}' and the singletino \tilde{S} with degenerate masses $\sim M_{Z'}$.

In the second example, we consider a special case in which the $U(1)'$ charge of H_2 is set

	Λ	M_Z		Λ	M_Z
h_u	0.88	1.0	h_d	0.06	0.09
h_e	0.04	0.04	h_D	0.70	0.85
h_s	0.36	0.31	A_s (GeV)	0	-48
A_u (GeV)	0	443	A_d (GeV)	0	555
A_e (GeV)	0	34	A_D (GeV)	0	494
m_1^2 (GeV) ²	(351) ²	(363) ²	m_2^2 (GeV) ²	(351) ²	-(815) ²
m_S^2 (GeV) ²	0	-(819) ²	$m_{D(c)}^2$ (GeV) ²	(1310) ²	(1440) ²

TABLE II. The initial and final values for the second example of the one singlet case. $n_D = 3$ and $\Lambda = 10^5$ GeV.

to be zero, i.e., $Q_1 = -Q_S = 5$. In this case, there is no $U(1)'$ D -term correction to the H_2 mass squared, reducing the value of $\tan\beta$ at the minimum of the potential. We present the numerical values of the parameters chosen for this example in Table II. Again, Λ is chosen to be 10^5 GeV and $n_D = 3$.

The vev of the singlet is $s = 3800$ GeV, the vev's of the Higgs doublets are $v_1 = 60$ GeV and $v_2 = 238$ GeV, resulting in a solution with $\mu = 818$ GeV and $\tan\beta = 4$, and thus $m_3^2 = (198 \text{ GeV})^2$. In this case, $M_{Z'} = 1140$ GeV and $\alpha_{Z-Z'} = 3.1 \times 10^{-4}$. The masses of the CP even physical Higgs are $m_{h_1} = 157$ GeV, $m_{h_2} = 434$ GeV, $m_{h_3} = 1130$ GeV at tree level, while $m_{h_1} = 181$ GeV at one loop. The CP odd Higgs scalar and the charged Higgs masses are $m_A = 409$ GeV and $m_{H^\pm} = 413$ GeV, respectively. The pattern of the spectrum is similar to that given in the previous example.

The masses of the charginos and neutralinos in this example are $m_{\tilde{\chi}_1^\pm} = 262$ GeV and $m_{\tilde{\chi}_2^\pm} = 843$ GeV; $m_{\tilde{\chi}_1^0} = 140$ GeV, $m_{\tilde{\chi}_2^0} = 262$ GeV, $m_{\tilde{\chi}_3^0} = 834$ GeV, $m_{\tilde{\chi}_4^0} = 841$ GeV, $m_{\tilde{\chi}_5^0} = 1130$ GeV and $m_{\tilde{\chi}_6^0} = 1140$ GeV. It again falls into the pattern of the chargino and neutralino spectrum discussed in the previous example.

In both examples squark and gluino masses are in the 1200 – 1400 GeV range. The next to lightest sparticle (NLSP) is the lightest neutralino, which is predominantly the bino, i.e., the gaugino of the $U(1)_Y$. The evolution of the parameters is illustrated in Figs. 1 and 2.

C. Multi-singlet models

We now consider the case with two singlets fields S and S' . Their charges are such that the coupling SS'^2 is allowed by gauge invariance. The F terms of the scalar potential (7) indicate that in addition to the effective μ term, $\mu = h_s\langle S \rangle$, there is an additional contribution to m_3^2 arising from the mixed term between $H_1 H_2$ and S'^2 . Hence, the total effective m_3^2 is given by

$$m_3^2 = h_s h_{s'} \langle S' \rangle^2 + A_s h_s \langle S \rangle, \quad (14)$$

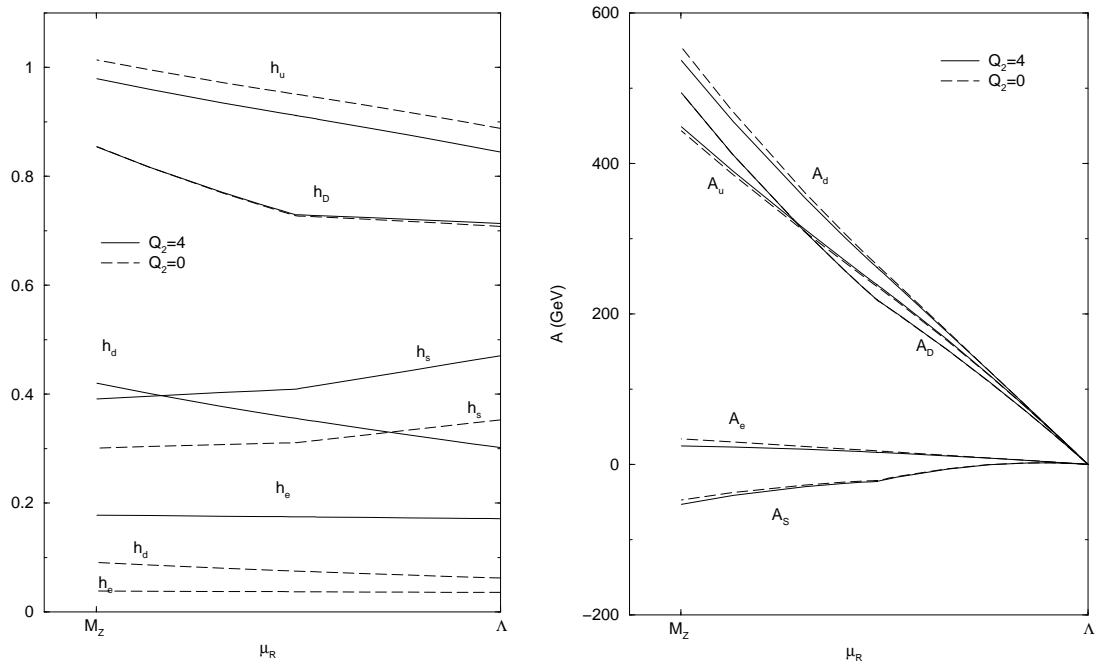


FIG. 1. The RGE evolution of the Yukawa and trilinear couplings in the two examples of the one singlet case. μ_R is the renormalization scale and $\Lambda = 10^5$ GeV.

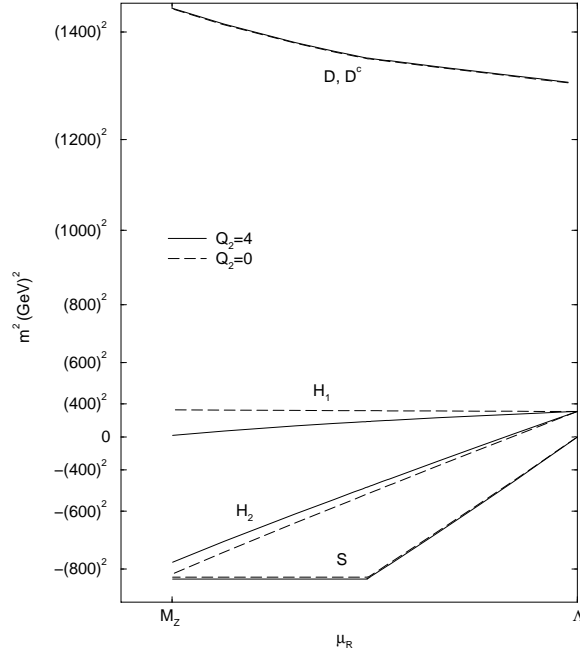


FIG. 2. The RGE evolution of the mass squared parameters in the two examples of the one singlet case.

	Λ	M_Z		Λ	M_Z
h_u	0.84	0.98	h_d	0.18	0.26
h_e	0.10	0.11	h_D	1.00	1.13
h_s	0.45	0.36	$h_{s'}$	0.12	0.10
A_s (GeV)	0	-50	$A_{s'}$ (GeV)	0	-29
A_u (GeV)	0	442	A_d (GeV)	0	538
A_e (GeV)	0	31	A_D (GeV)	0	477
m_1^2 (GeV) ²	(352) ²	(306) ²	m_2^2 (GeV) ²	(352) ²	-(758) ²
m_S^2 (GeV) ²	0	-(966) ²	$m_{S'}^2$ (GeV) ²	0	(15) ²
$m_{D(c)}^2$ (GeV) ²	(1280) ²	(1400) ²			

TABLE III. The initial and final values for the example of the two singlet case. $n_D = 3$ and $\Lambda = 10^5$ GeV.

when both S and S' acquire non-zero vev's. There are also corrections to the H_1 and H_2 mass-squared parameters from the D -term eqn. (8), $\delta m_{1/2}^2 = \frac{g_{1'}^2}{2} Q_{1/2} (Q_S s^2 + Q_{S'} s'^2)$, where $s^{(\prime)} = \sqrt{2} \langle S^{(\prime)} \rangle$ are the vev's of S and S' . The radiative $U(1)'$ symmetry breaking is again achieved by the coupling between S and the exotic quark pairs D and D^c . In this scenario, the self-coupling between S and S' stabilizes the vacuum of the two-singlet potential. To satisfy the phenomenological constraint on $M_{Z'}$, i.e., to ensure that the $U(1)'$ is broken at $\sim \mathcal{O}(\text{TeV})$, the Yukawa coupling $h_{s'}$ has to be small. Defining the vev's of the Higgs doublets as in the previous cases, the Higgs potential at the minimum again takes the form eqn. (13).

As an explicit example, we choose the charges of the Higgs doublets and the singlets to be $Q_1 = 1$, $Q_2 = 4$, $Q_S = -Q_1 - Q_2$ and $Q_{S'} = -Q_S/2$. We again choose $\Lambda = 10^5$ GeV and $n_D = 3$. We list the initial and final values of the relevant parameters in Table III.

The vev's of the singlets are $s = 4430$ GeV and $s' = 1760$, and those of the Higgs doublets are $v_1 = 23$ GeV and $v_2 = 245$ GeV, resulting in a solution with $\mu = 1110$ GeV and $\tan \beta = 10$, which is accidentally similar to the solution in our first example. The effective $m_3^2 = (407 \text{ GeV})^2$. The Z' mass is 1380 GeV, with the $Z - Z'$ mixing angle 0.004. The masses for the four CP even scalars, which are mixtures of the Higgs scalars and the scalar components from the singlets S and S' are $m_{h_1} = 124$ GeV, $m_{h_2} = 459$ GeV, $m_{h_3} = 1090$ GeV and $m_{h_4} = 1390$ GeV, with $m_{h_1} = 154$ GeV at one loop. The two CP odd scalar masses are $m_{A_1} = 138$ GeV and $m_{A_2} = 1080$ GeV. The charged Higgs scalar mass is $m_{H^\pm} = 1090$ GeV. Therefore, h_3 , A_2 and H^\pm approximately form an $SU(2)$ doublet which is not involved in the electroweak breaking. A_1 is predominantly associated with the two singlets. Its lightness is readily understood in terms of the small values of $h_{s'}$ and $A_{s'}$, and the extra global $U(1)$ symmetry which occurs for $h_{s'} = 0$ and $A_{s'} = 0$. h_4 is also mostly associated with the two singlets. As noted above, the μ parameter in this and the first example accidentally has similar values. However, the initial value of h_D has to be increased in this example (1 compared to 0.7 in the previous case) to counteract the effect of the self-coupling between the two singlets, so that the singlet vev's could both have large values.

The masses of the two charginos are $m_{\tilde{\chi}_1^\pm} = 266$ GeV and $m_{\tilde{\chi}_2^\pm} = 1130$ GeV; the masses of the seven neutralinos in this example are $m_{\tilde{\chi}_1^0} = 142$ GeV, $m_{\tilde{\chi}_2^0} = 266$ GeV, $m_{\tilde{\chi}_3^0} = 348$ GeV, $m_{\tilde{\chi}_4^0} = 1130$ GeV, $m_{\tilde{\chi}_5^0} = 1130$ GeV, $m_{\tilde{\chi}_6^0} = 1300$ GeV and $m_{\tilde{\chi}_7^0} = 1470$ GeV. It approximately follows the pattern of the chargino, neutralino spectrum discussed in the previous examples, which follows from limit of $v_1, v_2 \ll s(s')$ and the similar values of μ . However, $\tilde{\chi}_3^0$, coming from the self-coupling between the two singlet fields with mass $\sim h_{s'} s$, is a new feature. Squark and gluino masses are again in the 1200 – 1400 GeV range and the NLSP is again the lightest neutralino.

D. Variation of the parameters

As we discussed in Sec. II, the free parameters in our analysis are Λ , h_D and n_D . In our numerical examples, we choose $\Lambda = 10^5$ GeV. If Λ is varied while the other parameters are kept fixed, the scale at which the $U(1)'$ is broken changes so that Λ'/Λ is approximately a constant. For example, if we raise Λ to be 10^6 GeV in the first example, the singlet vev becomes $s \simeq 33400$ GeV. This is because RSB depends on the evolution interval, and not on the actual location of the boundary.

Next, consider n_D , the number of the exotic quarks that couple to the singlet field S with coefficient h_D . Choosing a smaller $n_D < 3$ and keeping Λ and h_D fixed, the $U(1)'$ breaking scale is reduced, and the associated Z' mass is diminished. For example, setting $n_D = 2$ in the first example instead of 3, the singlet vev is now $s = 3080$ GeV, resulting in $M_{Z'} = 919$ GeV with a mixing angle $\alpha_{Z-Z'} = 0.006$. In particular, with $n_D = 1$, $h_D = 0.7$ is too small to generate a solution with $M_{Z'} > 700$ GeV (which is a model dependent experimental lower bound). On the other hand, if n_D is fixed, increasing/decreasing h_D will raise/lower the $U(1)'$ scale and solutions for $n_D = 1$ exist for larger values of h_D . Hence, it is the product of h_D and n_D which is constrained by the Z' mass. Note that an upper bound on the Yukawa couplings $h_D \simeq \mathcal{O}(1)$ exists only if one requires the model to still be perturbative at some high energy scale. Also, the variation with n_D discussed above is similar for different exotic quark quantum numbers (for example, for two pairs of exotic quark singlets instead of an exotic quark doublet pair).

In our numerical analysis, we have identified, for simplicity, the messenger scale and the scale parameter Λ which determines the boundary conditions for the soft mass parameters eqs. (3) and (4). Differentiating these two scales will not change our conclusions. For example, raising the boundary scale for the RGE evolution in the first example to 5×10^5 GeV, while keeping $\Lambda = 10^5$ GeV in eqs. (3) and (4) leads to a solution with a larger singlet vev $s = 4430$ GeV due to the longer evolution interval. The A_s parameter is also larger in this case, $A_s = -86$ GeV, which results in a slightly smaller $\tan \beta = 13$ at the electroweak scale.

IV. SUMMARY AND CONCLUSIONS

It has been shown that a $U(1)'$ scale may be generated naturally and radiatively one or two orders of magnitude below the messenger scale. Upon integrating out the $U(1)'$ sector, the supersymmetry conserving (μ) and breaking (m_3^2) dimensionful Higgs mixing parameters are generated, resolving the μ problem in the otherwise attractive class of gauge-mediation models. A natural consequence of the models with only one singlet is that no new physical phases appear in the soft parameters, eliminating potentially unacceptable contributions to CP violating amplitudes which are flavor conserving and generically persist in gauge-mediation models. Other implications include new contributions to the effective quartic coupling which lift the lightest Higgs boson mass to $m_h \gtrsim 150$ GeV. The $U(1)'$ dynamics also adds to the already strong predictive power of the gauge-mediation framework, as the scalar, fermion and vector electroweak sectors are extended and new exotic matter is predicted at a few TeV scale. The Higgs, neutralino and chargino sectors are now extended and contain new degrees of freedom, but no new light pseudo Goldstone bosons. The additional gauge boson typically has a $\mathcal{O}(\text{TeV})$ mass and negligible mixing with the Z -boson. Its exact mass depends on the number of exotic quarks and on the strength of their Yukawa couplings.

We find the usual near equality between $|\mu|$ and the gluino mass that often appears in various variants of the MSSM. The gluino is heavy, which is a generic prediction of the gauge-mediation framework, and hence it naturally leads to a relatively large value of μ . At the same time it implies that electroweak symmetry breaking exhibits in our case the typical gauge mediation tuning $\sim |M_Z/M_{\text{gluino}}| \sim 1/10$, rather than a new tuning due to the $U(1)'$ dynamics. It is worth stressing that the Higgs mixing parameter in the scalar potential $m_3^2 = \mu A_s$ (in the one singlet case) is a geometrical mean of the superpotential Higgs mixing parameter μ and a radiatively generated (small) trilinear coupling A_s . Since μ is proportional to a large vev (and the heavy gluino implies further that $|\mu|$ is not suppressed by a small coupling) the geometrical mean $m_3 \sim \text{a few} \times 100$ GeV is sufficiently large.

While we have focused on the low-energy aspects of this one-scale model, it is interesting to consider its embedding in a high-energy theory, and in particular a unified or string theory. Unification of gauge couplings at Planckian energies constrains the exotic matter which is charged under the SM to fall into complete multiplets of a unified group, e.g., $D + D^c \rightarrow 5 + \bar{5}$ of $SU(5)$. It further constrains the number of such extra pairs of multiplets, e.g., $n_D = n_5 \lesssim 4$ [16]. The counting now includes also the messenger multiplets but may be modified by various considerations such as the exact mass (and hence, messenger) scale, the presence of large Yukawa couplings (as in our case) and the normalization of the $U(1)$ factor(s). Anomaly cancelation suggests that the SM, exotic multiplets (and messengers - if charged under $U(1)'$) are embedded in 27 multiplets of E_6 . The most straightforward realization which is consistent with the anomaly constraints (but which does not coincide with a standard E_6 model) is that only the third family carries $U(1)'$ charges, and the extra $5 + \bar{5}$ that are required for the anomaly cancelation correspond to one pair of exotics and to the usual Higgs doublets. The messengers are also not charged under the $U(1)'$ in this case. We have shown that such a model is consistent (for a large Yukawa coupling $h_D \sim 1$) but generically predicts a lighter Z' . This simple model, however, does not unify. Unified (and hence, automatically anomaly-free) models require a more complicated embedding,

typically contain more than one pair of exotic quarks, and further require some separation mechanism between the low-energy Higgs doublets and the heavier exotic leptons which appear. Such an embedding, however, is possible. An explicit example which satisfies all constraints (including unification of couplings) was given in Ref. [15].

Another issue that may arise is the kinetic mixing between $U(1)_Y$ and $U(1)'$. If at the fundamental scale, the mixed trace between $U(1)_Y$ and $U(1)'$ does not vanish, a term in the Lagrangian that mixes the field strength of the two $U(1)$'s could arise through loop corrections at the low energy scale. A non-orthogonal transformation of the gauge fields associated with the two $U(1)$'s is required so that the kinetic energy terms can be written in their canonical forms. This transformation effectively shifts the $U(1)'$ charges of the particles while keeping their $U(1)_Y$ charges [17]. The messengers in this case carry $U(1)'$ charges. It could also modify the low energy phenomenology. For example, the shifted charges for a specific model could be such that the low energy model is “leptophobic”, and hence allows a lighter Z' . However, the size of the correction depends in detail on the particle content of the model and on the decoupling scales. Kinetic mixing is again an ultra-violet effect (which, however, vanishes in the limit $g_{1'}/g_1 \rightarrow 0$) and was not considered here in detail.

Lastly, we would like to comment that a straightforward extension of our mechanism can lead to lepton number violation through couplings $h_\mu SLH_2 \rightarrow \mu_\mu LH_2$ if L and H_1 carry the same $U(1)'$ charge (this is the case for the $U(1)_\eta$ of E_6 , but not for the other $U(1)'$ embeddings) or more generally in a multi-singlet model. The case $Q_{H_1} \neq Q_L$ is in fact more attractive since it would forbid lepton number violating Yukawa operators in the high-energy theory. Since gauge mediation guarantees slepton-Higgs mass universality, and Higgs-slepton bilinear mixing in the scalar potential arises only from radiative A -parameters, then all conditions for the dynamical alignment suppression of neutrino masses outlined in Ref. [18] are automatically and naturally satisfied and gauge-mediation models would lead in this case to also a successful generation of neutrino masses.

In conclusion, the Higgs mass parameters in the gauge-mediation framework are best understood as dynamical degrees of freedom corresponding to a (SM) singlet. Here, it was suggested that such a singlet is not a gauge singlet but transforms under a $U(1)'$. While the minimal (“model-independent”) low energy framework was given and discussed, its many possible extensions and embeddings remain to be explored in greater detail.

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